

Ultra-cold WIMPs: relics of non-standard pre-BBN cosmologies

Graciela B. Gelmini¹ and Paolo Gondolo²

¹ Department of Physics and Astronomy, UCLA, 475 Portola Plaza, Los Angeles, CA 90095, USA

² Department of Physics, University of Utah, 115 S 1400 E # 201, Salt Lake City, UT 84112, USA

gelmini@physics.ucla.edu, paolo@physics.utah.edu

Weakly interacting massive particles (WIMPs) are one of very few probes of cosmology before Big Bang nucleosynthesis (BBN). We point out that in scenarios in which the Universe evolves in a non-standard manner during and after WIMP kinetic decoupling, the horizon mass scale at decoupling can be smaller and the dark matter WIMPs can be colder than in standard cosmology. This would lead to much smaller first objects in hierarchical structure formation. In low reheating temperature scenarios the effect may be large enough as to noticeably enhance indirect detection signals in GLAST and other detectors, by up to two orders of magnitude.

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In supersymmetric models, the lightest supersymmetric particle, usually a neutralino χ , is a good cold dark matter candidate. This is an example of a more general kind of cold dark matter candidates, weakly interacting massive particles (WIMPs). The relic density and velocity distribution of WIMPs before structure formation depend on the characteristics of the Universe (expansion rate, composition, etc.) before Big Bang Nucleosynthesis (BBN), i.e. at temperatures above $T \sim 1$ MeV. This is an epoch from which we have no data. Indeed, if dark matter (DM) WIMPs are ever found, they would be the first relics from that epoch that could be studied. Signatures of a non-standard pre-BBN cosmology that WIMPs may provide are few. Here we present one of, to our knowledge, only three.

In the standard scenario of WIMP decoupling, one assumes that the entropy of matter and radiation is conserved, that WIMPs are produced thermally, and, of most relevance for this work, that the first temperature of the radiation dominated epoch before BBN is high enough for WIMPs to have reached kinetic and chemical equilibrium before they decouple. WIMPs decouple first chemically and then kinetically. The chemical decoupling (or freeze-out) temperature T_{fo} is the temperature after which their number practically does not change, and in the standard case $T_{\text{fo-std}} \simeq m_\chi/20$, where m_χ is the WIMP mass. The kinetic decoupling temperature T_{kd} is the temperature after which WIMPs do not exchange momentum efficiently with the cosmic radiation fluid. Within the standard cosmology (SC), $T_{\text{kd-std}}$ lies between 10 MeV and a few GeV [1].

There are, however, well motivated cosmological models in which the standard assumptions above do not hold. These non-standard models include models with gravitino [2] or moduli [3] decay, Q-ball decay [4], thermal inflation [5], the Brans-Dicke-Jordan [6] cosmological model, models with anisotropic expansion [7] or quintessence domination [8]. It has been pointed out that in all of these models the neutralino relic density Ω_χ may differ from its standard value Ω_{std} (see e.g. Ref. [9]).

One clear signature of a non-standard cosmology before BBN would be WIMPs that compose part or all

of the DM but would be overabundant in the SC. Also the relic velocity distribution before structure formation in the Universe may differ from that in the SC. It has already been mentioned in the literature [10, 11] that WIMP's could be “hotter” than in the SC, even constituting warm instead of cold DM. This would leave an imprint on the large scale structure spectrum.

Here we point out a third possible signature of non-standard pre-BBN cosmologies: WIMPs may be “colder” (i.e. they may have smaller typical velocities and thus smaller free-streaming length) and the mass contained within the horizon at kinetic decoupling may be smaller than in the SC. This would lead to a smaller mass for the smallest WIMP structures, those formed first. Some of the smallest WIMPs clumps would survive to the present. Smaller and more abundant DM clumps would be present within our galaxy, an observable consequence of which would be a stronger annihilation signal from our galactic halo detected in indirect DM searches by GLAST, PAMELA and other experiments [12, 13, 14, 15, 16, 17, 18]. The signal in direct DM searches might also be affected in significant ways [19]. We show that in low reheating temperature cosmological models, both the free-streaming mass scale and the mass within the horizon at kinetic decoupling may be smaller than in the SC by many orders of magnitude

We present estimates of the kinetic decoupling temperature, characteristic relic WIMP velocity, free-streaming and kinetic decoupling horizon mass scales using a generic WIMP elastic scattering cross section written as in Ref. [14] and order of magnitude calculations. After chemical decoupling, $T \lesssim T_{\text{fo}}$, the total number of WIMPs remains constant and WIMPs are kept in local thermal equilibrium by elastic scattering with relativistic particles in the plasma. The WIMPs are non-relativistic at these temperatures, thus the average momentum exchanged per collision is small, of order T , and the rate of momentum exchange Γ is suppressed by a factor $\sim T/m_\chi$ with respect to the rate of elastic scattering

$$\Gamma \equiv \langle v \sigma_{\text{el}} \rangle n_{\text{rad}} \frac{T}{m_\chi} \simeq \sigma_0^{\text{el}} T^3 \left(\frac{T}{m_\chi} \right)^{2+l} . \quad (1)$$

Here σ_{el} is the total cross section for elastic scattering of WIMPs and relativistic Standard Model fermions, $n_{\text{rad}} \simeq T^3$ is the number density of relativistic particles, which are assumed to be in local thermal equilibrium, and $v \simeq 1$ is the WIMP-fermion relative velocity. The thermal average of σ_{el} can be written as $\langle \sigma_{\text{el}} \rangle = \sigma_0^{\text{el}} (T/m_\chi)^{1+l}$, where $\sigma_0^{\text{el}} \simeq (G_F m_W^2)^2 m^2 / m_Z^4 \simeq 10^{-10} m_\chi^2 \text{GeV}^{-4}$ sets the magnitude of the cross section, and l parametrizes its temperature dependence. Finally, m_W and m_Z are the masses of the standard gauge bosons and G_F is Fermi's coupling constant. In the Standard Model, elastic scattering between a light fermion and a heavy fermion is mediated by Z exchange and $l = 0$. In supersymmetric extensions of the Standard Model, where the lightest neutralino is the WIMP candidate, sfermion exchange occurs if the neutralino is a gaugino, Z exchange is suppressed, and $l = 1$. Inserting σ_0^{el} into Eq. 1, the rate of momentum exchange for non-relativistic WIMPs is

$$\Gamma \simeq \frac{10^{-10} T^5}{\text{GeV}^4} \left(\frac{T}{m_\chi} \right)^l. \quad (2)$$

In the following, we focus on the case of neutralinos with $l = 1$ (analogous results can be easily derived for $l = 0$).

Kinetic decoupling occurs when the rate of momentum exchange becomes smaller than the expansion rate of the Universe H . In the SC, decoupling happens while the universe is radiation dominated so the Hubble parameter is $H \simeq T^2/M_P$, where $M_P \simeq 10^{19} \text{GeV}$ is the Planck scale. From $\Gamma \simeq H$, we get

$$T_{\text{kd-std}} \simeq 20 \text{MeV} \left(\frac{m_\chi}{100 \text{GeV}} \right)^{1/4}. \quad (3)$$

More accurate calculations give a range of 10 MeV to a few GeV for $T_{\text{kd-std}}$ [1].

We concentrate now on a class of non-standard cosmological models with a late episode of inflation or entropy production [2, 3, 5, 9] in which a scalar field ϕ dominates the energy density of the Universe and subsequently decays (while oscillating around a minimum of its potential) eventually reheating the Universe to a low reheating temperature T_{RH} . This does not spoil primordial nucleosynthesis provided $T_{\text{RH}} \gtrsim 4 \text{ MeV}$ [20]. The interesting case for us is when T_{RH} is smaller than the standard chemical decoupling temperature $T_{\text{fo-std}}$, so kinetic decoupling happens during the ϕ -oscillations dominated phase.

Late-decaying scalar field models are well motivated in particle theories. For example, moduli fields which acquire a mass m_ϕ at the supersymmetry breaking scale 10 to 100 TeV and have gravitational strength interactions, thus their decay rate is $\Gamma_{\text{decay}} \simeq m_\phi^3/M_P^2$, are pervasive in supersymmetric models. These fields naturally tend to dominate the energy density of the Universe at late times and produce reheating temperatures in the MeV range (the “moduli problem” [21] is the tendency of these decays to happen even after BBN, which must be avoided). In fact, approximating the decay as instantaneous (usually a very good approximation) at the moment of decay

the energy stored in the field goes into radiation at a temperature T_{RH} . Thus $\Gamma_{\text{decay}} \simeq H(T_{\text{RH}}) \simeq T_{\text{RH}}^2/M_P$ implies that

$$T_{\text{RH}} \simeq 10 \text{ MeV} \left(\frac{m_\phi}{100 \text{ TeV}} \right)^{3/2}. \quad (4)$$

During the epoch in which the Universe is dominated by the oscillating ϕ field, the Hubble parameter H_ϕ is proportional to T^4 [22]. Since at the moment of ϕ decay, when $T = T_{\text{RH}}$, $H_\phi(T_{\text{RH}}) \simeq T_{\text{RH}}^2/M_P$, we can fix the proportionality constant so that $H_\phi \simeq T^4/(T_{\text{RH}}^2 M_P)$. Requiring that $\Gamma \simeq H_\phi$ at the new kinetic decoupling temperature $T_{\text{kd}'}$, we obtain

$$T_{\text{kd}'} \simeq 30 \text{ MeV} \left(\frac{10 \text{ MeV}}{T_{\text{RH}}} \right) \left(\frac{m_\chi}{100 \text{ GeV}} \right)^{1/2}. \quad (5)$$

Combining the two equations $\Gamma(T_{\text{kd-std}}) \simeq T_{\text{kd-std}}^2/M_P$ and $\Gamma(T_{\text{kd}'}) \simeq T_{\text{kd}'}^4/(T_{\text{RH}}^2 M_P)$, we have

$$T_{\text{kd}'} \simeq \frac{T_{\text{kd-std}}^2}{T_{\text{RH}}}. \quad (6)$$

Thus, if the reheating temperature is smaller than the standard kinetic decoupling temperature ($T_{\text{RH}} < T_{\text{kd-std}}$), WIMPs decouple earlier than in the SC, i.e. $T_{\text{kd}'} > T_{\text{kd-std}}$, and do so during the ϕ -oscillations dominated epoch.

So far we have assumed that the WIMPs are non-relativistic at decoupling. Thus, the relations above hold for $T_{\text{kd}'} < m_\chi/3$. If Eq. 6 leads to $T_{\text{kd}'} > m_\chi/3$, the WIMPs would be relativistic at decoupling, and the equations need to be modified. Since the momentum transfer to the radiation background is very efficient in collisions of relativistic WIMPs, in general the kinetic decoupling would happen at the moment WIMPs become non-relativistic and not earlier, i.e. $T_{\text{kd}'} \simeq m_\chi/3$ (unless the scattering cross section is so small that WIMPs are never in kinetic equilibrium [11]).

Late-decaying scalar field models exemplify many combinations of reheating and decoupling temperatures. In these models, the dominant WIMP production mechanism can be thermal (due to interactions with the radiation background) or non-thermal (due to the decay of the ϕ field into WIMPs), with or without chemical equilibrium (see for example Ref. [9]). In these models, neutralinos in almost all supersymmetric models could have the relic density necessary to be the DM [9] through a combination of thermal and non-thermal production mechanisms [3, 23, 24, 25, 26, 27, 28, 29, 30]. For thermal production without chemical equilibrium, most WIMP production happens at $T_* \simeq m_\chi/4$ [25]. Thus the WIMP number per comoving volume is fixed then. For $T_{\text{kd}'} < T_*$, the equations we derived above hold. For thermal production with chemical equilibrium the neutralino freezes out while the Universe is dominated by the ϕ field at a new freeze-out temperature $T_{\text{fo}'}$ higher than the usual $m_\chi/20$ [22, 26]. The freeze-out density is

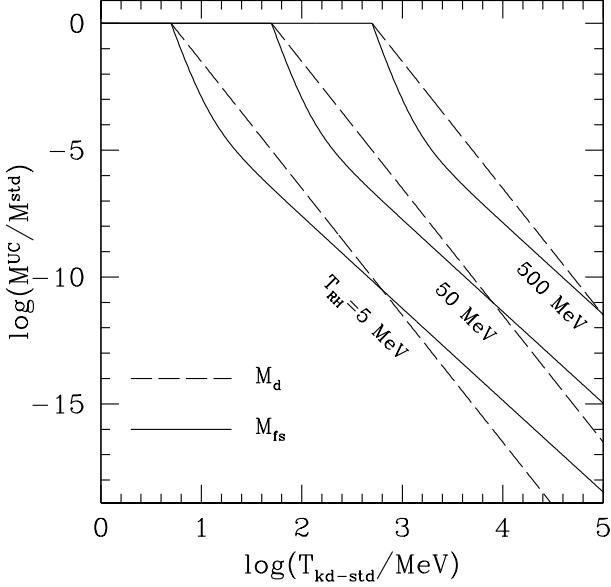


FIG. 1: Ratios of the late decaying scalar field (with ultra cold (UC) WIMPs) and standard (std) scenario free-streaming scales (M_{fs}^{UC}/M_{fs}^{std} , solid lines) and damping scales (M_d^{UC}/M_d^{std} , dashed lines), as functions of the standard kinetic decoupling temperature T_{kd-std} . Each set of two lines is labeled by the corresponding value of the reheating temperature T_{RH} .

larger than usual, but it is diluted by entropy production from ϕ decays, $\Omega_\chi \simeq T_{RH}^3 T_{fo-std}(T_{fo'})^{-4} \Omega_{std}$. The numerical results in Ref. [9] indicate a dependence of Ω_χ closer to T_{RH}^4 , and that the freeze-out temperature depends of T_{RH} as $T_{fo'}/20/m_\chi \simeq (T_{RH}20/m_\chi)^{1/4}$. In this case, the equations we derived hold for $T_{kd'} < T_{fo'}$. Non-thermal production without chemical equilibrium happens when production of WIMPs in the decay of the ϕ field is not compensated by annihilation. WIMPs are produced with an energy which is a fraction f of the ϕ -field mass m_ϕ , $E_\chi \simeq f m_\phi$. Thus if $m_\chi < 3 f m_\phi$, WIMPs are produced relativistic. In this case the elastic scattering cross section is $\langle v\sigma_{el} \rangle \simeq \sigma_0 T E_\chi / m_\chi^2 \simeq \sigma_0 T f m_\phi / m_\chi^2$. Taking the characteristic value of σ_0 as above, neutralinos are in kinetic equilibrium while they are relativistic for $m_\phi > 10 \text{ keV}/f(T_{RH}/10 \text{ MeV})^2$, that is for all the physically acceptable values of m_ϕ . Thus kinetic decoupling occurs while neutralinos are non-relativistic as assumed above. In the extreme case in which neutralinos are never in kinetic equilibrium [11], neutralinos can actually be warm DM [10], since they are produced late in the history of the Universe and with a large initial energy that redshifts until the moment of structure formation.

We now estimate the WIMP characteristic speed in low T_{RH} models. At the moment of kinetic decoupling, WIMPs are in thermal equilibrium with the radiation, and their characteristic speed is $v'(T_{kd'}) \simeq \sqrt{T_{kd'}/m_\chi}$. After decoupling, the speed decreases as a^{-1} (if the

WIMPs are non-relativistic). During the ϕ -oscillations dominated era, the scale factor of the Universe is related to the radiation temperature as $a \sim T^{-8/3}$. When the Universe becomes radiation dominated, i.e. at T_{RH} , the characteristic speed of WIMPs is therefore

$$v'(T_{RH}) \simeq \sqrt{T_{kd'}/m_\chi} (T_{RH}/T_{kd'})^{8/3}. \quad (7)$$

In comparison, in the standard radiation dominated case, $a \sim T^{-1}$ and at the same temperature T_{RH} WIMPs have a characteristic speed

$$v_{std}(T_{RH}) \simeq \sqrt{T_{kd-std}/m_\chi} (T_{RH}/T_{kd-std}). \quad (8)$$

Because speeds redshift in the same way in both models at temperatures smaller than T_{RH} , after reheating the speeds v' and v_{std} remain in the same ratio

$$v'/v_{std} \simeq (T_{RH}/T_{kd-std})^{10/3}. \quad (9)$$

This relation applies to the case $T_{kd'} < m_\chi/3$ for which Eq. 6 holds. Thus the characteristic relic WIMP speed in low T_{RH} cosmological models can be much smaller than in the SC. In other words, WIMPs can be much colder, i.e. “ultra-cold”, as we call them.

The free-streaming length λ_{fs} of ultra-cold WIMPs is consequently smaller than that of standard WIMPs. λ_{fs} is the characteristic distance covered by WIMPs from the time of kinetic decoupling t_{kd} to the present (while they propagate as free particles)

$$\lambda_{fs} = c a_0 \int_{t_{kd}}^{t_0} v \frac{dt}{a} \simeq c \sqrt{\frac{T_{kd}}{m_\chi}} a_0 a_{kd} \int_{a_{kd}}^{a_0} \frac{da}{a^3 H(a)}. \quad (10)$$

Here, T_{kd} and a_{kd} are the temperature and scale factor at the moment of kinetic decoupling, and a_0 is the value of the scale factor at present. During the ϕ -oscillations dominated epoch, $H(a) \propto a^{-3/2}$ with $a \sim t^{2/3} \sim T^{-8/3}$ and $H \simeq T^4/(T_{RH}^2 M_P)$. During the radiation dominated epoch, $H(a) \propto a^{-2}$, with $a \sim t^{1/2} \sim T^{-1}$ and $H \simeq T^2/M_P$, and the integral in the definition of λ_{fs} is $\propto \ln a$. During the matter-dominated epoch, $H(a) \propto a^{-3/2}$, with $a \sim t^{2/3} \sim T^{-1}$ and $H \simeq T^{3/2}/M_P$, and the free-streaming length saturates.

When the kinetic decoupling occurs during the ϕ -oscillations dominated epoch and WIMPs are ultra-cold (UC), we obtain

$$\lambda_{fs}^{UC} \simeq c \sqrt{\frac{T_{kd'}}{m_\chi}} \frac{a_{kd'} M_P}{a_0 T_0^2} \left\{ 2 \left[\left(\frac{T_{kd'}}{T_{RH}} \right)^{4/3} - 1 \right] + \ln \left(\frac{T_{RH}}{T_{eq}} \right) + 1.96 \right\}. \quad (11)$$

The first term within the curly brackets arises from the ϕ -dominated epoch, the logarithmic term from the radiation dominated epoch, and the term $1.96 = 2[1 - (1 + z_{eq})^{-1/2}]$ from the matter dominated epoch. The

subindex eq refers to matter-radiation equality. When Eq. 6 holds (for $T_{\text{kd}' < m_\chi/3}$), Eq. 11 becomes

$$\lambda_{\text{fs}}^{\text{UC}} \simeq \frac{cM_P}{T_0 \sqrt{m_\chi T_{\text{kd-std}}}} \left(\frac{T_{\text{RH}}}{T_{\text{kd-std}}} \right)^{23/6} \times \left\{ 2 \left[\left(\frac{T_{\text{kd-std}}}{T_{\text{RH}}} \right)^{8/3} - 1 \right] + \ln \left(\frac{T_{\text{RH}}}{T_{\text{eq}}} \right) + 1.96 \right\}. \quad (12)$$

This is to be compared with the free-streaming length in the SC,

$$\lambda_{\text{fs}}^{\text{std}} \simeq \frac{cM_P}{T_0 \sqrt{m_\chi T_{\text{kd-std}}}} \left[\ln \left(\frac{T_{\text{kd-std}}}{T_{\text{eq}}} \right) + 1.96 \right]. \quad (13)$$

As traditional, we introduce the mass M_{fs} contained within a sphere of radius $\lambda_{\text{fs}}/2$, and compare the standard and non-standard scenarios through the ratio

$$\frac{M_{\text{fs}}^{\text{UC}}}{M_{\text{fs}}^{\text{std}}} = \left(\frac{\lambda_{\text{fs}}^{\text{UC}}}{\lambda_{\text{fs}}^{\text{std}}} \right)^3. \quad (14)$$

As shown in Fig. 1 (solid lines), this ratio is always smaller than 1 if the reheating temperature is smaller than the standard kinetic decoupling temperature, and it can be many orders of magnitude smaller. For $T_{\text{RH}} = 5$ MeV and $T_{\text{kd-std}}$ between 10 MeV and a few GeV [1], the free-streaming scale M_{fs} can decrease by a factor between 0.1 and 10^{-13} (the suppression is less important for larger values of T_{RH}).

Friction between WIMPs and relativistic leptons during kinetic decoupling (Silk damping) leads to a small-scale cutoff in structure formation at the scale of the horizon at kinetic decoupling [15, 16]. The mass contained within the horizon at decoupling in the SC is $M_d^{\text{std}} \simeq 10^{-4} M_\odot (10 \text{ MeV}/T_{\text{kd-std}})^3$ [15, 16]. It varies from $10^{-4} M_\odot$ for $T_{\text{kd-std}} \simeq 10$ MeV to $10^{-12} M_\odot$ for $T_{\text{kd-std}} \simeq 5$ GeV.

During the ϕ -oscillations dominated phase the Universe expands, and thus the density contrast of DM inhomogeneities grow, in the same way as in a matter dominated phase. A detailed study (which is beyond the scope of this paper) of the kinetic decoupling during this phase should be done to find the cut-off mass scale of the smallest dark matter structures, which will also depend on the particular particle physics model considered. However, it is reasonable to assume that also in this case the cut-off will be given by the comoving free-streaming mass scale and/or the kinetic decoupling horizon mass scale.

During the ϕ -oscillations dominated phase the matter density scales as T^{-8} and the time as T^{-4} . Thus, the mass contained in the horizon at decoupling M_d^{UC} is smaller than in the SC by the factor

$$\frac{M_d^{\text{UC}}}{M_d^{\text{std}}} \simeq \left(\frac{T_{\text{RH}}}{T_{\text{kd}'}} \right)^4 \left(\frac{T_{\text{kd-std}}}{T_{\text{RH}}} \right)^3. \quad (15)$$

Using Eq. 6, this ratio becomes $(T_{\text{RH}}/T_{\text{kd-std}})^5$. As seen in Fig. 1 (dashed lines), the suppression factor

$(M_d^{\text{UC}}/M_d^{\text{std}})$ can be substantial, ranging from 0.1 for $T_{\text{kd-std}} = 10$ MeV to 10^{-15} for $T_{\text{kd-std}} = 5$ GeV, when assuming $T_{\text{RH}} = 5$ MeV (the suppression is less important for larger values of T_{RH}). This means that the range of M_d^{UC} is now from $10^{-27} M_\odot$ for $T_{\text{kd-std}} \simeq 5$ GeV to $10^{-5} M_\odot$ for $T_{\text{kd-std}} \simeq 10$ MeV.

From the ratios $(M_d^{\text{UC}}/M_d^{\text{std}})$ and $(M_{\text{fs}}^{\text{UC}}/M_{\text{fs}}^{\text{std}})$ shown in Fig. 1, we see that the damping mass scale M_d is less suppressed than the free-streaming mass scale M_{fs} , except possibly for standard kinetic decoupling temperatures above 1 GeV. This is important because the scale of the smallest WIMP haloes will be the larger of M_{fs} and M_d , since a halo mass must be larger than both. In the standard cosmological scenario, the damping scale M_d^{std} is usually larger than the free-streaming scale (by a factor $(m_\chi/T_{\text{kd-std}})^{3/2}$, at least for supersymmetric gaugino models with $T_{\text{kd-std}}$ in the 10–100 MeV range [15, 16]). Fig. 1 shows that in low T_{RH} models the damping scale remains larger than the free-streaming scale for most typical values of the standard kinetic decoupling temperature $T_{\text{kd-std}} \lesssim 1$ GeV. Thus, we take the damping scale M_d to be the characteristic mass of the smallest WIMP halos.

If the smallest halos survive until today, they may be present in the dark halo of our galaxy and may enhance the expected WIMP annihilation signals over the smooth halo expectation by a boost factor B . The B factor increases slowly with decreasing M_d . For a halo of mass M and smallest subhalo mass M_d , Ref. [31] finds $B \simeq 0.1[(M/M_d)^{0.13} - 1]$. In the SC, one expects B of the order of 10 for the Milky Way for which $M \simeq 10^{12} M_\odot$. For example, from the equation just mentioned we get $B \simeq 20$ to 130 for the standard range of M_d , from $10^{-6} M_\odot$ to $M_d \simeq 10^{-12} M_\odot$. With ultra-cold WIMPs, M_d could be much smaller and thus the boost factor could be much larger: it could reach $B \simeq 10^4$ for $M_d \simeq 10^{-27} M_\odot$. Such large boost factors would not only make a halo WIMP annihilation signal easier to detect, but would also be a signature of a non-standard pre-BBN cosmology.

We finally remark that ultra-cold WIMPs may arise in all models in which the expansion rate of the Universe in the pre-BBN era is larger than assumed in the SC, although the magnitude of the effect would in general be smaller than in the low reheating temperature models presented above. For example, let us consider “kination” models [8]. These are models in which the kinetic energy of a scalar field, $\rho_\phi = \dot{\phi}^2/2 \sim a^{-6}$, dominates the energy density of the Universe at $T > T_{\text{kin}}$ before BBN, while the entropy is dominated by the radiation, thus $a \sim T^{-1}$. Roughly, $\rho_\phi/\rho_{\text{rad}} \simeq \eta_\phi (T/\text{MeV})^2$, where η_ϕ is the fraction of the energy density of the Universe at $T \simeq 1$ MeV due to the kinetic energy of the ϕ field. At higher temperatures the fraction of ϕ kinetic energy grows very fast with the temperature and it is dominant for $T > T_{\text{kin}} \simeq \eta_\phi^{-1/2}$ MeV. If the kinetic decoupling of WIMPs happens during the kination period, i.e. if $T_{\text{kd-std}} > T_{\text{kin}}$, assuming $H \simeq \sqrt{\rho_\phi}/M_P$, the kinetic de-

coupling happens approximately at

$$T_{\text{kd}'}^{\text{kin}} \simeq 50 \text{MeV} \eta_{\phi}^{1/6} \left(\frac{m_{\chi}}{100 \text{GeV}} \right)^{1/3}. \quad (16)$$

In this case the free-streaming length is

$$\lambda_{\text{fs}}^{\text{kin}} \simeq c \sqrt{\frac{T_{\text{kd}'}}{m_{\chi}} \frac{a_{\text{kd}'} M_P}{a_0 T_0^2}} \left\{ 1 - \frac{T_{\text{kin}}}{T_{\text{kd}'}} + \ln \left(\frac{T_{\text{kin}}}{T_{\text{eq}}} \right) + 1.96 \right\}, \quad (17)$$

which for $T_{\text{kd}'}^{\text{kin}} > T_{\text{kd-std}} > T_{\text{kin}}$ is smaller than the standard free-streaming length. During the kination period, the scale factor is $a \sim T^{-1}$ and the time evolves as $t \sim T^{-3}$ thus the mass contained within the horizon at kinetic decoupling M_d^{kin} , again for $T_{\text{kd}'}^{\text{kin}} > T_{\text{kd-std}} > T_{\text{kin}}$, is smaller than the standard mass scale M_d^{std} by the ratio

$$\frac{M_d^{\text{kin}}}{M_d^{\text{std}}} \simeq \frac{T_{\text{kin}}^3 T_{\text{kd-std}}^3}{(T_{\text{kd}'}^{\text{kin}})^6}. \quad (18)$$

Thus in kination models the free-streaming and damping mass scales can be smaller than in the SC, although not by as much as in the late decaying scalar field (or low reheating temperature) models presented above.

We have here pointed out that a too large boost factor in the annihilation signal of a particular WIMP would be

a signature of a non standard cosmological evolution of the Universe just before BBN, during the kinetic decoupling of WIMPs. If dark matter WIMPs are ever found, they would be the first relics from the pre-BBN epoch that could be studied. Signatures of a non-standard pre-BBN cosmology that WIMPs may provide are few and here we presented one of them: WIMPs may be ultra cold so the mass of the smallest WIMP structures, those formed first, may be smaller than in the standard cosmology. Some of the smallest WIMPs clumps would survive to the present. Smaller and more abundant DM clumps would be present within our galaxy, an observable consequence of which would be a stronger annihilation signal from our galactic halo detected in indirect DM searches by GLAST, PAMELA and other experiments. Boost factors as large as 10^4 for usual WIMP candidates are possible in the low reheating temperature scenarios considered here. In last instance, verifying this signature would require to study in accelerators, the LHC or ILC, the properties of the particular WIMP that would allow us to estimate its scattering cross section, to find the same WIMP in indirect detection searches and to understand the formation and survival of the earliest dark matter clumps within the halo of our galaxy.

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[1] S. Profumo, K. Sigurdson and M. Kamionkowski, Phys. Rev. Lett. **97**, 031301 (2006).
[2] T. Moroi, M. Yamaguchi and T. Yanagida, Phys. Lett. **B 342**, 105 (1995); M. Kawasaki, T. Moroi and T. Yanagida, Phys. Lett. **B 370**, 52 (1996).
[3] T. Moroi and L. Randall, Nucl. Phys. **B 570**, 455 (2000).
[4] M. Fujii, K. Hamaguchi, Phys. Rev. D **66**, 083501 (2002); M. Fujii, M. Ibe, Phys. Rev. D **69**, 035006 (2004).
[5] D. H. Lyth, E.D. Stewart, Phys. Rev. D **53**, 1784 (1996).
[6] M. Kamionkowski and M. S. Turner, Phys. Rev. D **42**, 3310 (1990).
[7] J. D. Barrow, Nucl. Phys. B **208**, 501 (1982).
[8] P. Salati, Phys. Lett. B **571**, 121 (2003); S. Profumo and P. Ullio, JCAP **0311**, 006 (2003).
[9] G. B. Gelmini and P. Gondolo, Phys. Rev. D **74**, 023510 (2006); G. Gelmini, P. Gondolo, A. Soldatenko and C. E. Yaguna, Phys. Rev. D **74**, 083514 (2006).
[10] J. Hisano, K. Kohri and M. Nojiri, Phys. Lett. B **505**, 169 (2001).
[11] G. Gelmini and C. E. Yaguna, Phys. Lett. B **643**, 241 (2006).
[12] S. Hofmann, D. Schwarz and H. Stöcker, Phys. Rev. D **64** 083507 (2001); A. M. Green, S. Hofmann and D. Schwarz, Mon. Not. Roy. Astron. Soc. **353**, L23 (2004); J. Diemand, B. Moore and J. Stadel, Nature **433**, 389 (2005); B. Moore, J. Diemand, J. Stadel and T. Quinn, arXiv:astro-ph/0502213; H. Zhao, J. E. Taylor, J. Silk and D. Hooper, Astrophys. J. **654**, 697 (2007);
[13] T. Goerdt, O. Y. Gnedin, B. Moore, J. Diemand and J. Stadel, Mon. Not. Roy. Astron. Soc. **375**, 191 (2007); J. Diemand *et al.* arXiv:0805.1244 [astro-ph]; M. Kuhlen, J. Diemand and P. Madau, arXiv:0805.4416 [astro-ph].
[14] V. Berezinsky, V. Dokuchaev and Y. Eroshenko, Phys. Rev. D **68** 103003(2003),
[15] A. M. Green, S. Hofmann and D. J. Schwarz, JCAP **0508**, 003 (2005).
[16] A. Loeb and M. Zaldarriaga, Phys. Rev. D **71**, 103520 (2005).
[17] E. Bertschinger, Phys. Rev. D **74**, 063509 (2006).
[18] V. Berezinsky, V. Dokuchaev and Y. Eroshenko, Phys. Rev. D **73**, 063504 (2006);
[19] V. Berezinsky, V. Dokuchaev and Y. Eroshenko, and Phys. Rev. D **77**, 083519 (2008);
[20] M. Kamionkowski and S. M. Koushiappas, arXiv: 0801.3269 [astro-ph].
[21] M. Kawasaki, K. Kohri, and N. Sugiyama, Phys. Rev. Lett. **82**, 4168 (1999); Phys. Rev. D **62**, 023506 (2000); S. Hannestad, Phys. Rev. D **70**, 043506 (2004).
[22] G.D. Coughlan *et al* Phys. Lett. B **131**, 59 (1983); J.R. Ellis, D. V. Nanopoulos and M. Quiros, Phys. Lett. B **174**, 176 (1986); B. de Carlos *et al*, Phys. Lett. B **318**, 447 (1993); T. Banks, D. B. Kaplan and A. E. Nelson, Phys. Rev. D **49**, 779 (1994); S. Nakamura and M. Yamaguchi, Phys. Lett. B **665**, 167 (2007).
[23] J. McDonald, Phys. Rev. D **43** (1991) 1063.
[24] M. Kamionkowski, M. Turner, Phys. Rev. D **42** 3310

(1990); R. Jeannerot, X. Zhang, R. Brandenberger, JHEP **12**, 003 (1999); W. B. Lin, D. H. Huang, X. Zhang, R. Brandenberger, Phys. Rev. Lett. **86** 954 (2001).

[24] T. Moroi, M. Yamaguchi and T. Yanagida, Phys. Lett. **B 342**, 105 (1995); M. Kawasaki, T. Moroi and T. Yanagida, Phys. Lett. **B 370**, 52 (1996).

[25] D. J.H. Chung, E. W. Kolb and A. Riotto, Phys. Rev. D **60**, 063504 (1999).

[26] G. F. Giudice, E. W. Kolb and A. Riotto, Phys. Rev. D **64**, 023508 (2001).

[27] R. Allahverdi and M. Drees, Phys. Rev. Lett. **89**, 091302 (2002) and Phys. Rev. **D66**, 063513 (2002).

[28] S. Khalil, C. Muñoz and E. Torrente-Lujan, New Journal of Physics **4**, 27 (2002); E. Torrente-Lujan, hep-ph/0210036 (2002).

[29] N. Fornengo, A. Riotto, and S. Scopel, Phys. Rev. **D67**, 023514 (2003).

[30] C. Pallis, Astrop. Phys. **21**, 689 (2004).

[31] L. Strigari, S. Koushiappas, J. Bullock and M. Kaplinghat, Phys. Rev. D **75**, 083526 (2007).